

# Bulk Higgs Boson Decays in Brane Localized Gravity

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## Abstract

We embed the Standard Model in the Randall-Sundrum model of 5 dimensional brane localized gravity. The SM gauge and chiral fermion fields are restricted on the 4D visible brane whereas the Higgs and the right-handed neutrino are assumed to be 5D bulk fields. We calculate the effective couplings of the lowest mass Higgs field to the SM fermions and to the gauge bosons and find that the couplings are enhanced. Furthermore, the invisible decay width of a bulk Higgs of mass 150 GeV is shown to be large.

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It has been a long held belief in particle physics that the gravitational force is too feeble, i. e. the Planck scale too high, to have an impact on physics at the weak scale. Recent theoretical developments have shed new light on this sixteen order of magnitude disparity between the electroweak scale and the Planck scale, which is one of several hierarchy problems in particle physics. To this end, authors made ingenious uses of extra spatial dimensions which are present in many extensions of the Standard Model (SM) such as string theories and supergravity models. The number of extra dimensions  $n$  can range from 1 to 6 or 7. Another important ingredient is to confine the SM chiral fermions on one or more 3-branes which are stable topological objects in string theory. From the four dimensional field theory point of view these chiral fermions have only the usual Minkowski spacetime dependence given by  $x^\mu$  where  $\mu = 0, 1, 2, 3$  and not on the coordinates of the extra dimensions. On the other hand, gravity is allowed to spread into the extra dimensions. This opens up novel settings where new interplay between electroweak and gravitational physics can take place.

There are two main constructions to resolve the hierarchy problem under the general framework described above. The first one is to assume that the geometry of spacetime is factorizable and given by  $M^4 \times S^n$  where  $M^4$  denotes the usual four dimensional Minkowski space and the geometry of the extra dimensions is usually taken to be a  $n$ -tori for simplicity. The four dimensional Planck mass  $M_P$  is related to the fundamental scale of the higher dimensional theory  $M_*$  by the relation [1]

$$M_P^2 = M_*^{n+2} (2\pi R_1)(2\pi R_2) \dots (2\pi R_n) \quad (1)$$

where  $R_i$  with  $i = 1, 2, \dots, n$  denotes the compactification radii. For simplicity one takes all the radii to be equal to  $R$ . A very general prediction of this scenario is a modification of the Newtonian gravitational law which is well measured at large distances. From this one concludes that  $n \geq 2$  and the gravitational law is only changed at short distances below the micron range [1]. For  $n = 2$  astrophysical considerations set the limit on  $M_* \geq 50 \text{TeV}$  and  $R \leq 0.3 \mu\text{m}$  [2]. As  $n$  increases the constraints becomes less stringent.

An alternative scenario is given by Randall and Sundrum [3]. In the simplest version spacetime is taken to be five dimensional with the fifth dimension,  $y$ , compactified on a  $S_1/Z_2$  orbifold of radius  $r_c$ . Hence, we can write  $y = r_c \phi$  and  $-\pi \leq \phi \leq \pi$ . The points  $(x, \phi)$  and  $(x, -\phi)$  are identified. Two 3-branes with equal and opposite tensions are located at the orbifold fixed points: a visible brane at  $\phi = \pi$  where all the SM particles are confined and a hidden brane at  $\phi = 0$  where gravity is localized. The metric that solves the Einstein equations is given by [3]

$$ds^2 = e^{-2kr_c|\phi|} \eta_{\mu\nu} dx^\mu dx^\nu - r_c^2 d\phi^2 \quad (2)$$

where  $k$  is a parameter of the order of the fundamental mass scale  $M$  of the 5D theory and  $\eta_{\mu\nu}$  is the Minkowski flat metric with the signature  $(+---)$ . The exponential factor

in Eq.(2) is known as the warp or conformal factor. The 4D Planck scale is calculated to be given by

$$M_P^2 = \frac{M^3}{k} [1 - e^{-2kr_c\pi}] \quad (3)$$

and hence is of the order of the scale  $M$ . Furthermore, any field confined on the visible brane at  $\phi = \pi$  with mass  $m_0$  will be rescaled to have a physical mass given by  $m_0 e^{-kr_c\pi}$ . With the value of  $kr_c = 12$  a weak scale is dynamically generated with all fundamental masses of the order of  $M_P$ .

Treating Eq.(2) as a background metric, particles of different spins represented by full fledged bulk fields have been studied. The scalar field was treated in [4] and it was found that many of its properties are controlled by the warp factor. In addition it can be used to stabilize the extra dimension [5]. The SM singlet fermion is studied in [6] as a means of generating a small neutrino mass without using the seesaw mechanism. Issues of bulk gauge fields are examined in [7] and the embedding the SM in the full 5D bulk is given in [8], [9] and [10]. A characteristic of this scenario has emerged from these studies. In all cases the zero modes and the Kaluza-Klein (KK) excitations of the bulk fields are given by the roots of Bessel functions of different orders relating to their intrinsic spin. This is in sharp contrast to the case of factorizable geometry where the masses of the KK excitations are typically  $m^2 = \sum_i n_i^2 / R_i^2$  where  $n_i$  are integers. It was also noted in [8] that a fine tuning problem will emerge with the SM in the bulk if the Higgs boson is also allowed to extend into the fifth dimension.

In this paper we study a model with the SM chiral fermions and the gauge bosons all confined to a 3-brane at  $\phi = \pi$  but the Higgs boson is taken to be a bulk field. This is similar to the bulk scalar field studied in [4] but the bulk Higgs field develops a non-zero VEV after spontaneous symmetry breaking. The motivation here is to keep the feature that a bulk scalar field can confine fermions on a kink as found in [11]. Since we are confining the gauge fields on the brane, the Higgs field has local gauge symmetry on the visible brane. However, viewed in the fifth dimension, this is a global symmetry. Specifically, the brane fermions,  $\psi(x)$ , a  $U(1)$  brane gauge field  $A_\mu(x)$  and the bulk Higgs field,  $H$ , transform respectively as follow:

$$\psi(x) \rightarrow e^{i\Lambda(x)} \psi(x) \quad (4a)$$

$$A_\mu(x) \rightarrow A_\mu(x) + \partial_\mu \Lambda(x) \quad (4b)$$

$$H(x, \phi) \rightarrow e^{i\Lambda(x)} H(x, \phi). \quad (4c)$$

where the gauge function  $\Lambda$  depends on  $x^\mu$  only. We also include a bulk singlet fermion field denoted by  $\Psi(x^\mu, \phi)$  which will serve as a right-handed neutrino, like in [6]. The focus of our paper will be the decay modes of a physical Higgs of mass  $M_H$  between 125 and 250 GeV. With a mass in this range, the Higgs boson predominantly decays into a pair of b-quarks or gauge bosons. We shall see in detail later that this is also true for a

bulk Higgs boson. Hence we can examine quantitatively how the difference between the bulk Higgs and the SM Higgs boson can be manifested in the ongoing and future Higgs boson searches at high energy colliders.

The relevant action for the model described above can be written in four separate pieces. We begin with the bulk Higgs field action in 5D and it is given by:

$$S_H = \int d^4x \int_{-\pi}^{\pi} d\phi \sqrt{G} \left[ G^{AB} D_A H_0^\dagger D_B H_0 - \frac{\lambda}{4M} \left( H_0^\dagger H_0 - \frac{v_0^3}{2} \right)^2 \right] \quad (6)$$

where  $H_0$  denotes the bulk Higgs field, which is a weak doublet.  $G$  is the determinant of the metric tensor  $G_{AB}$  of Eq.(2).  $\lambda$  is a dimensionless parameter and  $v_0$  is the scale that characterizes the 5D vacuum expectation value of  $H_0$ . We use the notation that the capital Roman letters  $A, B, \dots$  denote 5D coordinates, the lower case letters  $a, b, \dots$  denote its tangent space coordinates, and the Greek letters  $\mu, \nu, \dots$  label Minkowski space coordinates.  $v_0$  is expected to be of order  $M$ . The gauge covariant derivative  $D_A$  is given by  $D_\mu = \partial_\mu + igA_\mu$  for the Minkowski coordinates where  $A_\mu$  is the 4D brane gauge field and  $D_5 = \partial_5$  for the fifth dimension.

We begin by studying the scalar sector. The real charge zero component of the doublet develops a non-zero VEV  $v_0^{\frac{3}{2}}/\sqrt{2}$ . After shifting the scalar field by  $H_0 \rightarrow (H + v_0^{\frac{3}{2}})/\sqrt{2}$ , we obtain an action similar to that given in [4]. To obtain the masses of the bulk Higgs boson and its KK excitations we first substitute the metric  $G_{AB}$  into Eq.(6) and use the Kaluza-Klein decomposition of  $H$ :

$$H(x, \phi) = \frac{1}{\sqrt{r_c}} \sum_n h_n(x) y_n(\phi). \quad (7)$$

Define that  $\sigma \equiv kr_c\phi$ . If the  $y_n(\phi)$  is chosen to satisfy the normalization condition:

$$\int_{-\pi}^{\pi} d\phi e^{-2\sigma} y_m(\phi) y_n(\phi) = \delta_{mn} \quad (8)$$

and the mass eigenvalue equation

$$-\frac{1}{r_c^2} \frac{d}{d\phi} \left( e^{-4\sigma} \frac{dy_n}{d\phi} \right) + m^2 e^{-4\sigma} y_n = m_n^2 e^{-2\sigma} y \quad (9)$$

with  $m^2 = \lambda v_0^3/M$ , the effective 4D action then simplifies to:

$$S_h^{(4)} = \frac{1}{2} \sum_n \int d^4x [\eta^{\mu\nu} \partial_\mu h_n \partial_\nu h_n - m_n^2 h_n^2] \quad (10)$$

which identify  $h_n$  as the  $n^{th}$  KK Higgs excitation with mass  $m_n$  given by Eq. (9). The lowest mass state will be the one we are interested in. It is useful to introduce the variables:

$$f_n \equiv e^{-2\sigma} y_n, \quad z_n \equiv \frac{m_n e^\sigma}{k}, \quad \text{and} \quad \omega \equiv \frac{m}{k} \quad (11)$$

Then Eq. (9) can be cast into the standard form:

$$z_n^2 \frac{d^2 f_n}{dz_n^2} + z_n \frac{df_n}{dz_n} + [z_n^2 - (4 + \omega^2)] f_n = 0. \quad (12)$$

The solutions of this equation are Bessel functions of order  $\nu = \sqrt{4 + \omega^2}$ :  $J_\nu(z_n)$ . It is clear that  $\nu \geq 2$  with the lower bound given by  $\lambda = 0$  which would be obtained by a higher dimension radiative symmetry breaking mechanism. After imposing the boundary conditions that the derivative of  $y_n$  be continuous at  $\phi = 0, \pi$  and the approximation  $e^{kr_c\pi} \gg 1$  and we obtain the following equation:

$$x_{n\nu} J_{\nu-1}(x_{n\nu}) = (\nu - 2) J_\nu(x_{n\nu}) \quad (13)$$

which gives the eigenvalues of  $x_{n\nu} \equiv \frac{m_n}{k} e^{kr_c\pi}$ . For the case of  $\omega \simeq 0$  the eigenvalues are essentially the roots of  $J_1(x)$ . Numerically the first two values are 3.83 and 7.02. The lowest eigenvalues of  $x_1$  vary between 3.8 to 12.5 when  $\omega = 0.1 \sim 10$ . We identify the lowest mode to be the lightest Higgs boson of mass  $m_1$ ; then

$$k\epsilon \equiv k e^{-kr_c\pi} = \frac{m_1}{x_1}. \quad (14)$$

In Table 1 we give  $k\epsilon$  for different values of  $\omega$  and  $m_1$ . The masses of the first KK excited states are also shown. It can be seen that in this scenario we are led to a tower of closely spaced Higgs states for a wide range of  $\omega$  values.

Now we continue with the gauge boson-Higgs interaction contained in Eq.(6). Using standard notations the action for the bulk Higgs- $W$  interaction is given by

$$\int d^4x \int_{-\pi}^{\pi} d\phi \sqrt{G} [G^{\mu\nu} W_\mu^- H_0^\dagger W_\nu^+ H_0] \quad (15)$$

After spontaneous symmetry breaking, the  $W$  boson mass is generated via:

$$\frac{1}{4} \int d^4x r_c e^{-4kr_c\pi} e^{2kr_c\pi} g^2 [\eta^{\mu\nu} W_\mu^- W_\nu^+ v_0^3] \quad (16)$$

From this we can relate the physical  $W$  mass,  $M_W$ , to the 5D parameters; i.e.

$$M_W = \frac{1}{2} g v_0 e^{-kr_c\pi} \sqrt{r_c v_0} \quad (17)$$

or equivalently we can write the Fermi scale  $v$  as

$$v = \epsilon v_0 \sqrt{r_c v_0} = 250 \text{ GeV}. \quad (18)$$

We can now calculate the coupling of the KK Higgs to  $W^+W^-$ . Using Eqs.(15),(7), (11) and noting that the bulk scalar fields are evaluated at the visible brane, we find the  $h_n W^+ W^-$  coupling is

$$g M_W \sqrt{\frac{kr_c}{1 - (\frac{\omega}{x_{n\nu}})^2}} \quad (19)$$

Compare this with the case in SM where the coupling is  $gM_W$ . Similar conclusion applies to the Higgs- $Z$  boson coupling. Hence, in this model the ratio of the width of the lowest Higgs state decaying into 2 gauge bosons compared to that of the SM Higgs decay is

$$R_g = \frac{\Gamma(h_1 \rightarrow W^+W^-, ZZ)}{\Gamma(H_{SM} \rightarrow W^+W^-, ZZ)} = \frac{kr_c}{1 - (\frac{\omega}{x_{1\nu}})^2}. \quad (20)$$

We note that  $x_1$  is determined by the order of the Bessel equation and hence on  $\omega$  but not on the choice of  $m_1$ ; then  $R_g$  takes the values 12 and 33.3 for  $\omega = 0.1, 10$  respectively and it is most sensitive to the choice of  $kr_c$ .

Next we introduce the brane fermions located at  $\phi = \pi$  and their interactions with  $H$ . Our discussions will be given in terms of the SM lepton doublet  $L_0$  and the right-handed  $e_{0R}$  and the subscript 0 is used to denote unrenormalized fields. The results obtained can be easily carried over to the quarks. Using the notation that  $g_{\mu\nu}^{vis}(x_\mu) = G_{\mu\nu}(x_\mu, \phi = \pi)$  and  $g^{vis}$  its determinant, the action is given by

$$S_{bf} = \int d^4x \sqrt{-g^{vis}} \left( \bar{L}_0 \hat{\gamma}^\mu \partial_\mu L_0 + \bar{e}_{0R} \hat{\gamma}^\mu e_{0R} \right) - \frac{\hat{y}_e}{\sqrt{M}} \int d^4x \sqrt{-g^{vis}} \bar{L}_0 H_0 e_{0R} + h.c. \quad (21)$$

where

$$\begin{aligned} \hat{\gamma}^\mu &= E_a^\mu(\phi = \pi) \gamma^a \\ &= e^{kr_c \pi} \gamma^\mu. \end{aligned} \quad (22)$$

In Eq.(21)  $H_0$  is at  $\phi = \pi$  and in Eq.(22) the inverse vielbein,  $E_a^A = \text{diag}(e^\sigma, e^\sigma, e^\sigma, e^\sigma, \frac{1}{r_c})$ . A ubiquitous Yukawa coupling,  $\hat{y}_e$ , is introduced in Eq.(21) and is a free parameter. The SM chiral fermions reside on the brane. The field  $H$  will be expanded via Eq.(7) and evaluated at  $\phi = \pi$ . The kinetic term will require a rescaling due to nontrivial  $g^{vis}$  and  $E_a^A$ . The brane fermion wavefunction rescaling is

$$L_0 = e^{\frac{3}{2}kr_c\pi} L \quad (23)$$

and the kinetic term for  $L$  is in the canonical form. The  $e_{0R}$  field is similarly rescaled. Spontaneous symmetry breaking will generate a mass for the electron via the Yukawa term in Eq.(21). This is given by

$$m_e = \hat{y}_e \epsilon v_0 \sqrt{\frac{v_0}{2M}} \quad (24)$$

A second product of this manipulation is an expression for the effective Yukawa coupling of the interaction  $\bar{e}_L e_R h_n$  on the visible brane and it is

$$y_e^{eff} = \frac{gm_e}{2M_W} \sqrt{\frac{kr_c}{1 - (\frac{\omega}{x_{n\nu}})^2}} \quad (25)$$

where we have used Eqs.(11), (13) and (17). Barring fortuitous tuning of parameters it is clear that the decay width of the lowest state bulk Higgs,  $h_1$ , into a fermion pair will be different from the SM Higgs boson due to the square root factor in Eq.(25). Indeed the ratio of the width of  $h_1$  decaying to a fermion pair to that of a SM Higgs of the same mass is predicted to be

$$R_f = \frac{\Gamma(h_1 \rightarrow \bar{f}f)}{\Gamma(H_{SM} \rightarrow \bar{f}f)} = \frac{kr_c}{1 - (\frac{\omega}{x_{1\nu}})^2} \quad (26)$$

Noticed that the gauge boson width and all the fermion widths are enhanced by the same factor:  $R_f = R_g$ .

This is a convenient point to pause and take stock of the parameters of the model. Aside from the Yukawa coupling there are five 5D parameters, *viz.*  $v_0, M, k, r_c$ , and  $m$ , which replaces  $\lambda$ . As argued in [3] that  $kr_c$  is of O(10) in order to solve the hierarchy problem. Without loss of generality we take  $kr_c = 12$  as a benchmark value. We have also seen that  $\omega = m/k$  is of order unity and this is a parameter we vary as was done in Table 1. Then by choosing a value for  $m_1$  the value of  $k$  and  $v_0$  is determined by Eq.(13) and (18). Similarly the fundamental 5D scale  $M$  will be fixed by Eq.(3).

As noted in [6] that a SM singlet bulk fermion,  $\Psi$ , can be added to the brane localized gravity model and generate a small Dirac neutrino mass for  $\nu_{eL}$  as an alternative to the seesaw mechanism. The construction of the action for  $\Psi$  is given in [6] and we generalized it to include a bulk Higgs mechanism. The action of the bulk fermion takes the form [12]

$$S_\Psi = \int d^4x \int_{-\pi}^{\pi} d\phi \sqrt{G} \left\{ E_a^A \left[ \frac{i}{2} \bar{\Psi} \gamma^a (\partial_A - \overleftarrow{\partial}_A) \Psi + \frac{\omega_{bcA}}{8} \bar{\Psi} \{ \gamma^a, \sigma^{bc} \} \Psi \right] - m_D \text{sgn}(\phi) \bar{\Psi} \Psi \right\} \quad (27)$$

where we have included a bare Dirac mass  $m_D$  and  $\omega_{bcA}$  is the spin connection associated with the warp metric and it gives no contribution to the physics we are discussing. A convenient choice of the Dirac matrices is  $\gamma^a = (\gamma^\mu, i\gamma^5)$ . In 5D one can define the left- and right-handed spinors  $\Psi_{L,R} = \frac{1 \mp \gamma^5}{2} \Psi$ . We are mainly interested in  $\Psi_R$  and will present enough formulas so that we are self contained and our notations are clear. Details of their derivations can be obtained in [6]. After using the KK decompositions of  $\Psi_{L,R}$  is

$$\Psi_{L,R} = \frac{1}{\sqrt{r_c}} \sum_n e^{2\sigma} \psi_{nL,R}(x) \hat{f}_{nL,R}(\phi) \quad (28)$$

The functions  $\hat{f}_{nL,R}$  are separate complete sets of orthonormal functions. After imposing the  $\mathbb{Z}_2$  orbifold symmetry and the periodic boundary condition of  $\Psi_{L,R}(x, \pi) = \Psi_{L,R}(x, -\pi)$  the range of  $\phi$  is the interval  $[0, \pi]$ , and we get the 4D action for the bulk neutrino and its KK excitation

$$S_\psi = \sum_n \int d^4x (\bar{\psi}_n i \gamma^\mu \partial_\mu \psi_n - \mu_n \bar{\psi}_n \psi_n) \quad (29)$$

where  $\psi = \psi_L + \psi_R$  and  $\mu_n \geq 0$ . As in the scalar case the functions  $\hat{f}_{nL,R}$  satisfy

$$\int_0^\pi d\phi e^\sigma \hat{f}_{nR}^* \hat{f}_{mR} = \int_0^\pi d\phi e^\sigma \hat{f}_{nL}^* \hat{f}_{mL} = \delta_{nm} \quad (30)$$

and the equation

$$\left( \pm \frac{1}{r_c} \partial_\phi - m_D \right) \hat{f}_{nL,R} = -\mu_n e^\sigma \hat{f}_{nR,L}. \quad (31)$$

With the following change of variables:

$$t \equiv \epsilon e^\sigma, \quad \tilde{\nu} \equiv \frac{m_D}{k}, \quad \text{and} \quad \tilde{x}_n \equiv \frac{\mu_n}{\epsilon k} \quad (32)$$

and

$$f_{nL,R}(t) = \frac{\hat{f}_{nL,R}(\phi)}{\sqrt{k r_c \epsilon}}, \quad (32a)$$

$$\tilde{f}_{nL,R}(t) = \frac{\hat{f}_{nL,R}(t)}{\sqrt{t}} \quad (32b)$$

we can combine the two first order equations into the standard Bessel equation of half integer order:

$$t_n^2 \frac{d^2 \tilde{f}_{nL,R}}{dt_n^2} + t_n \frac{d \tilde{f}_{nL,R}}{dt_n} + \left[ t_n^2 - (\tilde{\nu} \mp \frac{1}{2}) \right] \tilde{f}_{nL,R} = 0 \quad (33)$$

where  $t_n \equiv \tilde{x}_n t$ . In particular we are interested in the zero mode with  $\mu_n = 0$  which has the normalization

$$|f_{0R}(1)|^2 = \frac{1 - 2\tilde{\nu}}{1 - \epsilon^{1-2\tilde{\nu}}}. \quad (34)$$

For  $\tilde{\nu} \geq \frac{1}{2}$  this suppresses the wavefunction at the visible brane by the factor  $\epsilon^{\tilde{\nu}-\frac{1}{2}}$ . On the other hand the higher KK bulk neutrino modes are not subjected to such a suppression. Similar to the bulk Higgs the KK bulk neutrinos have masses given by the equation

$$J_{\tilde{\nu}-\frac{1}{2}}(\tilde{x}_n) = 0 \quad (35)$$

which can be seen [Eq.(32)] to be of order weak scale. Thus we expect no more than one or two such neutrino modes are kinematically accessible to light Higgs boson decays. This is strikingly different from bulk neutrinos in the factorizable geometry scenario [13] where a large number of neutrinos are available if the radius of compactification is sufficiently large. The phenomenology of the simplest model of this type is discussed in [14] and [13].

We now proceed to discuss the interaction between the brane leptons, the bulk Higgs bosons and the bulk neutrinos. The action is:

$$S_{b\nu} = -\sqrt{2} \hat{Y}_5 \int d^4x \sqrt{g_{vis}} \bar{L}_0 H_0(x, \pi) \Psi_R(x, \pi) + h.c. \quad (36)$$



where  $\hat{Y}_5$  is yet another Yukawa coupling. After substituting in the KK decompositions of Eqs.(7), (28), and Eqs. (33) plus fields rescaling and spontaneous symmetry breaking we arrive at the couplings between  $\nu_L$  and the KK bulk neutrino states. These will provide the off diagonal terms for neutrino mass matrix. Explicitly, they are given by

$$\sum_n \hat{Y}_5 \int d^4x \epsilon v_0 r_c \sqrt{v_0 k} f_{nR}(\pi) \bar{\nu}_L \psi_{nR} + h.c. \quad (37)$$

For the bulk zero mode we obtain a mass term for  $\nu_L$  with the value given by

$$m_\nu = \hat{Y}_5 v \sqrt{kr_c} \epsilon^{\tilde{\nu}-\frac{1}{2}} \quad (38)$$

where we have used Eqs.(34) and (18). We shall demand that  $10^{-4} \text{ eV} \leq m_\nu \leq 2 \text{ eV}$  as indicated by current solar and atmospheric neutrino data [15] as well as direct measurement of neutrino mass in tritium  $\beta$  decays [16]. Assuming that  $\hat{Y}_5$  is of order 1 and using the values we have obtained previously on the parameters in Eq.(38) we find that  $\tilde{\nu}$  lies in the range 1.1 to 1.5 corresponding to 2 eV and  $10^{-4} \text{ eV}$  light neutrino respectively. As can be seen this conclusion is not sensitive to  $\hat{Y}_5$  unless it is extremely large or small. For the lower value of  $\tilde{\nu} \sim 1.1$  the roots of the half integer Bessel function is close to  $j\pi$  where  $j$  is an integer. Thus, we find that diagonal terms of neutrino mass matrix of the KK neutrinos is well approximated by

$$\mu_n \simeq n\pi k \epsilon \quad (39)$$

Similarly for  $\tilde{\nu} \sim 1.5$  the solutions of Bessel functions of order 1 apply. Both are independent of  $\hat{Y}_5$ . It is interesting to note that the small value of  $m_\nu$  is a result of the function  $f_{1R}$  being almost vanishing at  $\phi = \pi$  [6] which in turn is due to the warp metric and the assumed separation between the hidden and the visible branes. For details of the neutrino mass matrix see [6].

By an analogous calculation we derive the effective coupling between  $\nu_L$ , the  $m^{th}$  bulk neutrino mode, and the  $n^{th}$  KK Higgs boson which we denote by  $\tilde{y}_{nm}^{eff}$ . Thus,

$$\begin{aligned} \tilde{y}_{nm}^{eff} &= \hat{Y}_5 y_n(\pi) f_{mR}(\pi) \sqrt{kr_c} e^{-kr_c \pi} \\ &= \hat{Y}_5 kr_c \sqrt{\frac{2}{1 - (\frac{\omega}{x_{n\nu}})^2}} \end{aligned} \quad (40)$$

where we used the normalized values of  $y_n$  and  $f_{nR}$ . We note that this effective coupling is not small and universal for each of the KK Higgs excitations; however, it does vary from one Higgs excitation to another even though numerically this variation is not large. To get an idea of the numerics we take the lightest Higgs boson,  $h_1$ , and  $\omega=1$ . This results in  $\tilde{y}_1^{eff} = 17.5\hat{Y}_5$  and  $\tilde{y}_2^{eff} = 17.1\hat{Y}_5$ . The enhancement over the naive Yukawa coupling is due to the fact that the wavefunctions of neither the bulk Higgs nor the neutrino states are small on the visible brane. Without resorting to fine tuning of the Yukawa coupling

this leads to a large invisible width of the  $h_1$  if the channel is kinematically open. For example, the width of  $h_1$  into  $\bar{\nu}_L\psi_{1R}$  is given by

$$\Gamma_{\nu\psi_1} = \Gamma(h_1 \rightarrow \bar{\nu}_L\psi_{1R} + \bar{\psi}_{1R}\nu_L) = \frac{m_1}{8\pi} \frac{(\hat{Y}_5 k r_c)^2}{1 - \frac{\omega^2}{x_{1\nu}^2}} \left(1 - \frac{\mu_1^2}{m_1^2}\right)^2 \tan^2 \theta_\nu, \quad (41)$$

where  $\tan \theta_\nu$  is the mixing between the lightest neutrino and its KK excitation. In principle this can be constrained by the invisible width of the Z boson if the the KK neutrino is lighter than 90 GeV. For heavier neutrinos we have to rely on details of the model, instead we shall take this to be a free parameter. For  $\hat{Y}_5 = \frac{1}{3}$ ,  $\omega = 1$  and  $\tan \theta_\nu \sim 0.1$ , a 150 GeV Higgs boson will have an invisible width,  $\Gamma_{inv} = 170$  MeV. This is to be compared with the  $b\bar{b}$  width,  $\Gamma_b$  which we calculated to be 32.8 MeV using Eq.(25) and the SM width,  $\Gamma_b^{SM} = 2.6$  MeV. We have also used the running  $b$ -quark mass which is the dominant QCD correction. We also note the following interesting branching ratio:

$$\frac{\Gamma_{\nu\psi_1}}{\Gamma_b} = \frac{1}{3\sqrt{2}} \frac{k r_c}{G_F} \left(\frac{\hat{Y}_5}{m_b}\right)^2 \left(1 - \frac{\mu_1^2}{m_1^2}\right)^2 \left(1 - \frac{4m_b^2}{m_1^2}\right)^{-\frac{3}{2}} \tan^2 \theta_\nu \quad (42)$$

In Fig. 1 we display the branching ratio into  $b\bar{b}$  of  $h_1$  for the case with bulk neutrinos as a function of the mass  $m_1$ . The parameters used are given previously. The corresponding SM Higgs branching ratio is also given as a comparison. One striking feature is the invisible width which is large for the values of  $\hat{Y}_5$  and  $\tan \theta_\nu$  we used and this suppresses the branching ratio. In Fig.2 the total decay width of the  $h_1$  is plotted for different values of  $m_1$ . The enhancement factors are evident.

In conclusion we have constructed a model of bulk Higgs boson in the Randall-Sundrum scenario where the SM chiral fermions and gauge bosons are confined on the visible brane. By identifying the first KK excitation of the 5D scalar as the lowest mass Higgs we find that the partial widths and hence the total decay width of such a boson are enhanced compared to the SM. This enhancement is directly proportional to the quantity  $k r_c$ . This model also predicts that the effective Higgs to  $t\bar{t}$  is large which does not necessarily mean that perturbation is not applicable since the Yukawa coupling itself can still be of order unity. A detail investigation of this issue is certainly worthwhile but is beyond the scope of the present paper.

Furthermore, the phenomenology of the bulk Higgs changes drastically if we add bulk neutrinos into the picture. This has the added motivation of providing a mechanism for generating a small mass for the active  $\nu_L$ . In this case, the invisible width of the Higgs boson is not negligible even though only one or two KK bulk neutrino decays are open for Higgs masses between 125 to 250 GeV. Obviously, this has important consequences for Higgs boson searches in high energy colliders. Other characteristics of the model include the mass spectra of the Higgs and bulk neutrino which are given by roots of Bessel functions of various orders. Surprisingly, for the bulk neutrinos for the phenomenologically interesting cases the order of the Bessel function varies between 1/2 to 1. This makes the model much more predictive than naively expected. We eagerly await a rich phenomenology of the model discernibly different from the SM to be discovered.

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Table 1: The mass of the first KK excitation,  $m_2$  and values of  $k\epsilon$  in parenthesis for different values of  $\omega$  (first column) and  $m_1$  in GeV (first row)

$\omega \backslash m_1$	125	150	200	250
0.1	228 (32.6)	275 (39.1)	366 (52.1)	458 (65.2)
1.0	224 (30.6)	266 (36.7)	358 (48.9)	448 (61.1)
2.0	213 (26.4)	255 (31.7)	340 (42.3)	425 (52.8)
10.0	169 (10.0)	203 (12.0)	270 (16.0)	338 (20.0)

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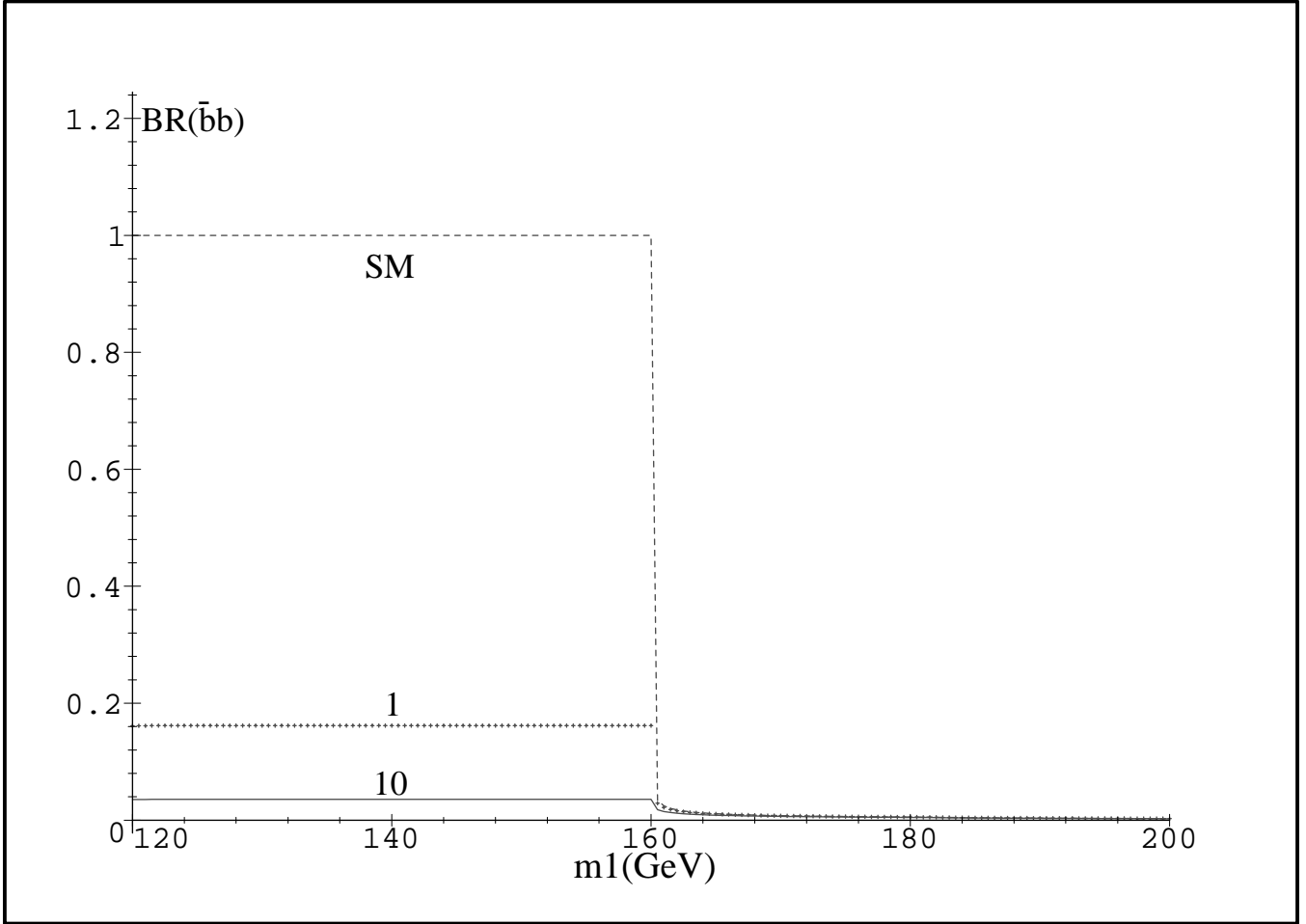


Figure 1: The branching of the lowest mass Higgs into  $b\bar{b}$  as a function of the mass  $m_1$ . The values for  $\omega = 1$  and 10 are shown. The SM Higgs result is given by the dashed curve.

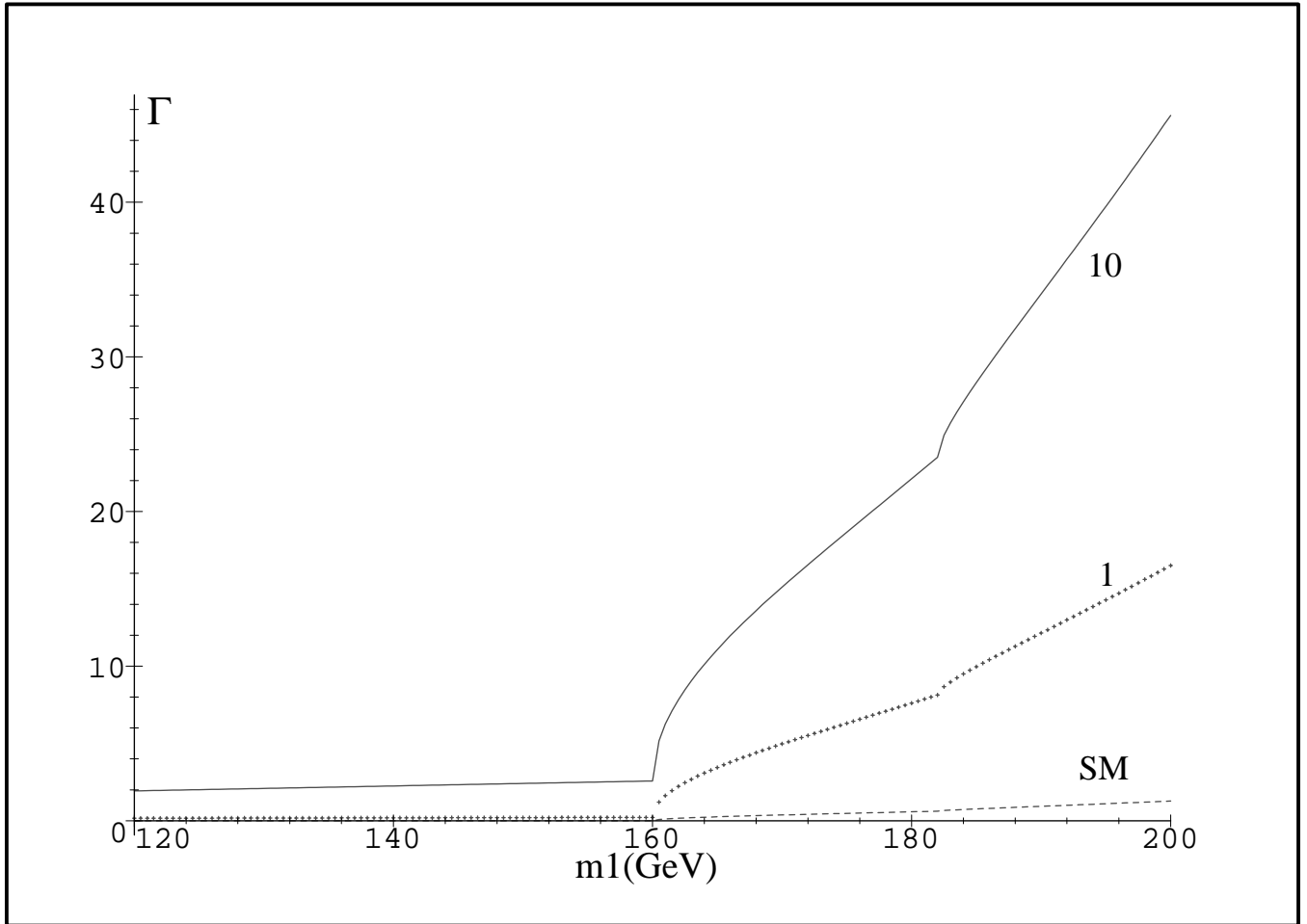


Figure 2: The total decay width of  $h_1$  as a function of  $m_1$ . The width for the SM Higgs boson is given by the dashed line